Heliosheath fluctuations near the perpendicular termination shock: two-dimensional hybrid simulations

Kaijun Liu, S. Peter Gary, and Dan Winske

Los Alamos National Laboratory, Los Alamos, New Mexico

Kaijun Liu, S. Peter Gary, and Dan Winske, Los Alamos National Laboratory, Los Alamos, NM 87545. (e-mail: kaijun@lanl.gov, pgary@lanl.gov, winske@lanl.gov)

Abstract.   The two-dimensional Los Alamos hybrid simulation code is used to study the excitation of fluctuations and the associated ion dynamics at the high Alfvén-Mach heliospheric termination shock and in the near-shock heliosheath. This simulation represents the electrons as a zero-mass fluid, addresses only a perpendicular shock, and considers the upstream ions to consist of a cool solar wind component and a more energetic pickup component. The shock yields two strongly heated downstream components: the more dense thermals and the more tenuous suprathermals, each with strong $T_\perp > T_\parallel$ anisotropies. In the downstream, relatively homogeneous sheath plasma, linear dispersion theory predicts that each component anisotropy drives both the Alfvén-cyclotron instability and the proton mirror instability. The simulation demonstrates that Alfvén-cyclotron modes dominate mirror-like modes in the downstream, in contrast to earlier two-dimensional hybrid simulations of high-Mach quasi-perpendicular shocks without pickup ions, in which the Alfvén-cyclotron and mirror modes were excited to comparable intensities. A hypothesis is presented that the presence of pickup ions implies a relatively low magnetosonic Mach number at termination shocks, thereby favoring the excitation of the Alfvén-cyclotron instability. The primary consequence of wave-particle interactions from this instability is pitch-angle scattering, so the energetic part of the ion perpendicular velocity distribution in the heliosheath is depleted by comparison with the distribution computed from a comparable one-dimensional hybrid simulation.

1. Introduction

As the solar wind flows away from the Sun, interstellar neutral particles, primarily hydrogen, become ionized, are picked up and become part of the flow. So in the outer heliosphere the flowing plasma primarily consists of a more dense, relatively cool ($T \sim$ eV) solar wind ion component, and a more tenuous, more energetic ($\sim$ keV) pickup ion component. The decreasing energy density of this outward flowing plasma leads to a transition from supersonic to subsonic flow; this is the termination shock where the flowing plasma is compressed and heated.

The Voyager 1 spacecraft crossed the termination shock in 2004 [*Stone et al.*, 2005; *Burlaga et al.*, 2005] and Voyager 2 made a number of crossings of the shock in late 2007 [*Stone et al.*, 2008; *Burlaga et al.*, 2008]. The Voyager 2 crossings were notable in that low-energy (< 10 eV) plasma measurements, as well as magnetic field measurements, were available, so that heating of the thermal ion component was measured. The magnetic field profiles, especially those of the TS-2 crossing of Voyager 2, showed a well-defined shock structure, including an extended rise of the magnetic field ("foot") upstream, a sharp rise of the magnetic field ("ramp") at the shock itself, and a peak of the magnetic field at the end of the ramp ("overshoot") that is larger than the average downstream value. Such features are also characteristic of quasi-perpendicular bow shock crossings for upstream Alfvén Mach numbers larger than a critical value (a "supercritical shock") [e.g. *Greenstadt et al.*, 1980; *Bale*, 2005].

However, neither Voyager spacecraft directly measured the properties of the pickup ions, and so we have no *in situ* observations of the response of these ions to the termination

shock and its environs. Ion kinetic simulations of the termination shock reproduce the

qualitative observed features of the shock crossing [*Liewer et al.*, 1993; *Lee et al.*, 2005; *Wu et al.*, 2009], indicating how the pickup ions modify the basic spatial scales of the shock

and how plasma flow energy dissipated at the shock is partitioned between the pickup

and solar wind ions. But we do not have, and are not likely to soon obtain, detailed

measurements of species velocity distributions at and near this shock.

Beyond the termination shock lies the heliosheath, where the plasma is further processed

by wave-particle scattering and by charge-exchange with the interstellar neutrals. Quasi-

perpendicular shocks drive ion heating preferentially perpendicular to the background

magnetic field; this excites instabilities which, via wave-particle scattering, reduces the

ion anisotropy. Charge-exchange between the heated ions and cold interstellar neutrals

leads to an enhanced flux of energetic neutral particles, predominantly hydrogen, which

propagate away from this region unaffected by the electromagnetic fields of the interplan-

etary and interstellar medium.

The Interstellar Boundary Explorer, or IBEX [*McComas et al.* 2004], is an Earth-

orbiting satellite with the primary objective of studying the global interaction between

the solar wind and the interstellar medium. IBEX carries two highly sensitive Energetic

Neutral Atom (ENA) imaging instruments which measure neutral particles in two energy

bands, 10 eV to 2 keV and 300 eV to 6 keV. In look directions away from the vicinity of

the Earth, the primary source of such neutrals is the inner heliosheath; by imaging these

neutral particles, IBEX is providing new information about how our heliosphere interacts

with the interstellar medium which surrounds it [*McComas et al.*, 2009; *Funsten et al.*,

2009]. Heating at the termination shock and subsequent scattering in the heliosheath

71 determine the velocity distributions of the protons which in turn determine the fluxes of

72 the ENAs measured by IBEX.

73 Downstream of the quasi-perpendicular terrestrial bow shock, both observations [*An-*

74 *derson et al.*, 1994; *Farrugia et al.*, 2005] and simulations [*Thomas and Brecht*, 1986,

75 1987; *Winske and Quest*, 1988; *McKean et al.*, 1995, 1996] have demonstrated strong ion

76 anisotropies of the type $T_\perp/T_\parallel > 1$, leading to enhanced field fluctuations from both the

77 Alfvén-cyclotron and proton mirror instabilities. Observations of the magnetic field in

78 the heliosheath have been made from both Voyager 1 and 2 [*Burlaga et al.*, 2005; 2008],

79 showing large amplitude compressive fluctuations called magnetic holes [*Burlaga et al.*,

80 2007]. But to date there are no reported measurements of thermal ion anisotropies or cor-

81 responding instability-driven enhanced fluctuations in the heliosheath. Alfvén-cyclotron

82 and mirror-like fluctuations scatter ions differently, so it is important to determine their

83 relative amplitudes in the heliosheath if we are to predict the properties of ions heated by

84 the termination shock and the consequent ENA fluxes which emanate from this region.

85 Just as the upstream ions consist of two distinct components, the more dense, cooler

86 solar wind (denoted by subscript sw) and the more tenuous, hotter pickup ions (denoted by

87 subscript pu), termination shock simulations show that the heated downstream ions also

88 consist of two distinct components [*Wu et al.*, 2010]. In our treatment of linear theory as

89 applied to downstream instabilities described in the Appendix, we follow *Wu et al.* [2010]

90 and use the adjective "thermal" (subscript t) to denote the more dense, relatively cool

91 ions and "suprathermal" (subscript s) to denote the more tenuous, hotter ions. However,

92 in our simulations we can label and then follow each upstream ion and distinguish solar

93 wind from pickup components in the downstream. Relatively few solar wind ions are

94 accelerated by the shock up to suprathermal energies, so in Section 4 and Section 5 we

95 continue to use the terms "solar wind" and "pickup" to refer to the respectively cooler and

96 hotter downstream ion components, with the understanding that these are approximately

97 equivalent to the corresponding labels of thermal and suprathermal.

98 This manuscript describes two-dimensional hybrid simulations of the perpendicular ter-

99 mination shock with emphasis on the excitation of enhanced fluctuations at and down-

100 stream of the shock. Our particular interest here is how these fluctuations interact with

101 the few KeV suprathermal ions which are a primary source of ENA fluxes measured

102 by IBEX instrumentation. Section 2 summarizes relevant results from earlier simulations.

103 Section 3 reviews properties of the Los Alamos hybrid simulation code, Section 4 describes

104 general plasma results from the simulations, and Section 5 discusses the fluctuating fields

105 resulting from the computations. Section 6 presents a summary and a conclusion, while

106 Appendix A presents results of linear theory for instabilities downstream of the simulated

107 shocks.

108 We denote the ion plasma frequency as $\omega_i \equiv \sqrt{4\pi n_o e^2 / m_i}$, the ion cyclotron frequency

109 as $\Omega_i \equiv e_i B_o / m_i c$, the ion inertial length $\lambda_i \equiv c/\omega_i$, and the β of the jth ion component

110 as $\beta_{\|j} \equiv 8\pi n_j k_B T_{\|j} / B_o^2$. The Alfvén speed is $v_A = B_o / \sqrt{4\pi n_o m_i}$. Here n_o is the total

111 upstream plasma density and $\mathbf{B}_o$ denotes the upstream uniform background magnetic

112 field. The Cartesian coordinate systems of our two-dimensional simulations and of linear

113 dispersion theory are the same; there are spatial variations in both the direction parallel

114 to $\mathbf{B}_o$ (denoted by $\|$) and one direction perpendicular to the background field (denoted

115 by $\perp$), but no spatial variations in the other perpendicular direction (denoted by $\perp\perp$).

Here θ denotes the angle between $\mathbf{k}$ and $\mathbf{B}_o$: $\hat{\mathbf{k}} \cdot \hat{\mathbf{B}}_o = \cos(\theta)$. In the simulation, $\hat{\mathbf{x}}$ is the $\perp$ direction, $\hat{\mathbf{y}}$ is the $\parallel$ direction, and $\hat{\mathbf{z}}$ is the $\perp\perp$ direction.

2. Previous Simulations

The most powerful and most complete tools for understanding the plasma physics of collisionless shocks are kinetic simulations. There are two distinct categories of such computations: particle-in-cell (PIC) and hybrid simulations.

PIC simulations provide a self-consistent description of the evolution of plasma ions and electrons as well as the associated electric and magnetic fields. Although one-dimensional PIC simulations have been applied to the termination shock [*Lee et al.* 2005; *Chapman et al.*, 2005; *Matsukiyo et al.*, 2007], such computations require that the electron dynamics be followed in detail. This restricts such calculations to relatively short times and, typically, to one spatial dimension.

Hybrid simulations treat ion species as superparticles, but represent electrons as a fluid while providing a self-consistent treatment of the electric and magnetic fields. Hybrid codes can run to much longer times and/or in more spatial dimensions than PIC codes, and thus are well suited for studying ion responses to space plasma disturbances such as the termination shock. The trade-offs are that hybrid codes cannot describe kinetic electron physics and represent only fluctuations with frequencies near or below the proton cyclotron frequency and with wavelengths near or longer than the ion inertial length. But this is precisely the regime of strong interactions between the waves and the ions, so our hybrid code is an appropriate tool for studies of ion dynamics in collisionless plasmas.

Liewer et al. [1993] used a one-dimensional hybrid code to show that for high Mach number quasi-perpendicular termination shocks, the primary mechanism for pickup ion

138 energy gain is due to those particles which backstream from downstream to upstream

139 where they gain energy before passing back downstream through the shock. As the shock

140 normal angle becomes less than about 60°, oblique and quasi-parallel shocks develop.

141 For such shocks pickup ion heating becomes more involved; reflected ions propagate back

142 upstream, exciting magnetosonic waves which further scatter and energize the ions [*Liewer*

143 *et al.* 1993, 1995; *Kucharek and Scholer*, 1995; *Scholer et al.* 2002].

144 Recently *Wu et al.* [2009; 2010] used the one-dimensional Los Alamos hybrid code to

145 simulate the perpendicular termination shock. *Wu et al.* [2009] developed a modified

146 Rankine-Hugoniot model that included pickup ions to interpret the hybrid simulations.

147 *Wu et al.* [2010] showed that specular reflection of ions [*Zank et al.*, 1996] by the electric

148 field normal to such a shock is not necessary for an initial energy gain of pickup ions;

149 rather, the large gyroradius of these ions enables some of them to return upstream where

150 they gain substantial energy from the motional electric field as they gyrate in the "foot"

151 [*Lee et al.*, 2005; *Chapman et al.*, 2005; *Matsukiyo et al.*, 2007]. Further, *Wu et al.* [2010]

152 demonstrated that faster solar wind flows lead to increased fluxes of ions in the tails

153 of their velocity distributions, consistent with energetic neutral atom observations from

154 IBEX [*Funsten et al.*, 2009].

155 Almost all kinetic termination shock simulations published to date have considered spa-

156 tial variations only in one dimension. Such simulations can produce genuine shocks that

157 satisfy the Rankine-Hugoniot relations (e.g., *Wu et al.*, 2009). However, one-dimensional

158 simulations have the fundamental limitation that downstream heating occurs only in the

159 two directions perpendicular to the background magnetic field. Two-dimensional sim-

160 ulations of the quasi-perpendicular bow shock without pickup ions have demonstrated

downstream ion perpendicular temperature anisotropies and associated Alfvén-cyclotron and mirror-like enhanced fluctuations [*Thomas and Brecht*, 1986, 1987; *Winske and Quest*, 1988; *McKean et al.*, 1995, 1996] that scatter the ions and provide heating in all three dimensions. Fully three-dimensional simulations of the perpendicular bow shock (*Thomas*, 1989) show very similar shock structure compared to two-dimensional runs, with slight modifications due to the further relaxing of constraints on the allowable modes. And recent two-dimensional and three-dimensional simulations of the competition between Alfvén-cyclotron and proton mirror instabilities in a uniform plasma modeling the region downstream of the bow shock show differences at late times, as large nonlinear three-dimensional structures can develop from mirror-mode fluctuations [*Shoji et al.*, 2009].

In contrast to the many one-dimensional kinetic simulations, *Giacalone and Decker* [2010] used two-dimensional hybrid simulations of the termination shock with an imposed spectrum of upstream turbulence to address ion acceleration to MeV energies due to the combination of these imposed fluctuations and those which develop self-consistently. In contrast, the two-dimensional simulations of the termination shock described here do not admit an initial spectrum of enhanced fluctuations, but rather investigate in detail the field fluctuations which arise self-consistently at the shock or by Alfvén-cyclotron and proton mirror instability growth downstream.

The Alfvén-cyclotron instability has maximum growth rate at $\mathbf{k} \times \mathbf{B}_o = 0$ and under that condition, it yields only fluctuating magnetic field components in the plane perpendicular to $\mathbf{B}_o$. Thus enhanced fluctuations from this instability have $k_\perp \simeq 0$ and $\delta B_\parallel \simeq 0$. The proton mirror instability has maximum growth rate at propagation oblique to $\mathbf{B}_o$, so that $k_\perp \neq 0$ and $k_\parallel \neq 0$, but linear theory demonstrates that the fluctuating magnetic

184 field components associated with this growing mode lie predominantly in the $\|$-$\perp$ plane;

185 thus for mirror-like fluctuations $\delta B_{\perp\perp} \simeq 0$.

186   The primary wave-particle interaction from enhanced fluctuations from the Alfvén-

187 cyclotron instability is through the ion cyclotron resonance.   This implies that

188 the principal consequence of this interaction is pitch-angle scattering of ions

189 near or above the ion thermal speed.   In contrast, enhanced fluctuations

190 from the proton mirror instability interact mostly through the Landau resonance

191 on subthermal ions.   So, Alfvén-cyclotron wave scattering acts largely on the

192 pickup(upstream)/suprathermal(downstream) ions, whereas mirror mode scattering has

193 its greatest effect on the solar wind(upstream)/thermal(downstream) ion components. To

194 understand the different consequences of these two instabilities we need physical insight

195 into the downstream competition between the two.   A prime purpose of the research

196 described here is to provide such insight.

197   The competition between Alfvén-cyclotron and mirror-like fluctuations downstream of

198 quasi-perpendicular shocks is sometimes expressed in terms of growth rates computed from

199 linear theory [*Gary*, 1992] and simulations of the instabilities in homogeneous plasmas

200 [*McKean et al.*, 1992; 1994; *Shoji et al.*, 2010]. In this scenario, ion anisotropies driven at

201 the shock are the source of linear instability growth as the plasma flows downstream into

202 the relatively homogeneous sheath, as in the low Alfvén-Mach bow shock simulations of

203 *McKean et al.* [1996]. However, the high Alfvén-Mach bow shock simulations of *McKean*

204 *et al.* [1995] provided a different picture, in which nonlinear processes immediately at

205 the inhomogeneous shock drive large-amplitude fluctuations which decay as the plasma

206 is convected downstream, where linear theory may be used to identify the modes. The

₂₀₇ *McKean et al.* [1995] simulations of the quasi-perpendicular shock, which do not include

₂₀₈ pickup ions, show that the Alfvén-cyclotron and proton mirror modes are present at

₂₀₉ comparable intensities downstream.

3. Los Alamos hybrid simulation code

₂₁₀ The two-dimensional Los Alamos hybrid code [*Winske and Omidi*, 1993; 1996; *Winske*

₂₁₁ *et al.*, 2003] self-consistently computes the evolution of plasma ions, fluid electrons, and

₂₁₂ the associated fluctuating electric and magnetic fields. This is a robust, well-documented

₂₁₃ simulation code that has been extensively used to study both fast and slow shocks in space

₂₁₄ plasmas and other environments [*Winske and Quest*, 1988; *Yin et al.*, 2005; *Winske and*

₂₁₅ *Gary*, 2007]. In this code plasma ions are represented as superparticles, while electrons are

₂₁₆ treated as a massless, adiabatic fluid with $\gamma = 5/3$. The model allows spatial variations in

₂₁₇ both the direction of shock propagation perpendicular to $\mathbf{B}_o$ (the x direction) and in the

₂₁₈ direction parallel to $\mathbf{B}_o$ (the y direction). The x- and y- dimensions of the computational

₂₁₉ box are $L_x = 400\lambda_i$ and $L_y = 256\lambda_i$, the number of cells in the x- and y dimensions are

₂₂₀ $N_x = 800$ and $N_y = 256$, and the computational time step $\Omega_i\Delta t = 0.02$.

₂₂₁ The simulation is run in the downstream, or heliosheath, plasma frame, so that the

₂₂₂ incident ions are injected from the left wall with an average right-directed velocity; as

₂₂₃ these ions approach the right-hand wall, the shock is formed and propagates back to the

₂₂₄ left. As in *Wu et al.* [2009; 2010], our simulation is of a strictly perpendicular shock;

₂₂₅ our initial conditions are that the upstream solar wind flow speed v_{sw} corresponds to

₂₂₆ an Alfvén Mach number $M_A = v_{sw}/v_A = 8$, an upstream solar wind $\beta_{sw} = 0.04$, and

₂₂₇ an upstream pickup ion relative density of $n_{pu}/n_o = 0.20$. The upstream pickup ion

₂₂₈ velocity distribution is assumed to be the isotropic (in the solar wind frame) Vasyliunas-

229 Siscoe distribution [*Vasyliunas and Siscoe*, 1976] modified as in *Wu et al.* [2010] so that

230 $f_{pu}(v) \sim 1/v^{3/2}$; here $\beta_{pu} = 8.53$.

4. Simulation: General results

231 This section describes some general results from the simulation of the perpendicular

232 termination shock with parameters as described in Section 3. Figure 1 uses black lines to

233 illustrate the profiles of the main component of the magnetic field $B_y(x)$ at $\Omega_i t = 60$ and

234 at three different values of y given in the figure caption. The shock is at $x \simeq 197\lambda_i$ at this

235 moment. This figure shows substantial variability in the magnetic profile as a function of

236 y, just as Voyager 2 observed a variety of different profiles during its multiple crossings of

237 the termination shock [*Burlaga et al.*, 2008]. Nevertheless, the y-averaged magnetic field,

238 indicated by the gray line in Figure 1 and denoted as $\overline{B_y}$, exhibits the main magnetic

239 features of a perpendicular shock, namely, the foot, ramp, and overshoot, as well as

240 downstream oscillations along x, similar to the magnetic fluctuations computed in one-

241 dimensional simulations [*Wu et al.*, 2009, 2010]. These downstream fluctuations decrease

242 with increasing x faster than in the comparable one-dimensional shock simulations. This

243 is consistent with the time-stationary (with respect to the shock) oscillations observed at

244 a clean perpendicular shock at Earth by Cluster [*Horbury et al.*, 2001; *Bale*, 2005] and is

245 associated with the relaxation of the downstream plasma anisotropies, which is precluded

246 in the one-dimensional model.

247 Figure 2 displays y-averaged plasma properties of the solar wind ion component (left

248 panels) and the pickup ions (right panels) as functions of x at $\Omega_i t = 60$. Panels from

249 top to bottom present the jth ion component densities averaged over y, the component

250 temperatures, the component temperature ratios $T_{\perp j}/T_{\|j} = (T_{xj} + T_{zj})/(T_{yj}/2)$, and the

component $\beta_{\|j}$, respectively. The temperatures here are calculated from $T_{x,y,z} = m_i <$

$v_{x,y,z}^2 - < v_{x,y,z} >^2 >$, where the notation "$<>$" means the velocity-averaged value of the

enclosed quantity over all the ions of a particular component. Note that the solar wind

ions and the pickup ions are not necessarily Maxwellian at places, especially in the vicinity

of the shock.

Consistent with the one-dimensional results of *Wu et al.* [2009], the temperature increase

of the solar wind ions across the shock (~ 50) is much larger than that of the pickup ions,

but the energy gain of the pickup ions accounts for most ($> 80\%$) of the shock dissipation.

Accordingly, $\beta_{\|sw}$ shown in Figure 2g jumps from its upstream value of 0.04 to about 0.8

across the shock associated with the significant increase of T_{swy} displayed in Figure 2c.

The solar wind ions have a large anisotropy near the shock as shown by the big spike

(~ 25) of $T_\perp/T_\|$ of the solar wind ions in Figure 2e. This anisotropy declines to ~ 1.5

immediately downstream of the shock. Across the shock, the average $\beta_{\|pu}$ does not change

significantly except for some oscillations associated with the changes in $\overline{B_y}$ as shown in

Figure 1. The pickup ion $T_\perp/T_\|$ increases to 2.5 near the shock and gradually declines to

about 1.2 far downstream.

Using the plasma parameters illustrated in Figure 2, we have solved the kinetic lin-

ear dispersion equation for two locations: immediately downstream of the shock near

$x/\lambda_i \simeq 197$ and somewhat further downstream near $x/\lambda_i \simeq 210$. At both locations,

the Alfvén-cyclotron instabilities have appreciably faster growth rates than their proton

miorror instability counterparts. Immediately downstream, theory predicts that the very

strong anisotropy of the thermal component is the dominant instability driver, whereas

273 futher downstream, after $T_{\perp t}/T_{\|t}$ is substantially reduced, the weaker suprathermal ion

274 anisotropy becomes the primary driver of the Alfvén-cyclotron enhanced fluctuations.

5. Simulations: Fluctuating fields

275 In addition to the Alfvén-cyclotron and mirror-like fluctuations discussed in Section 2,

276 previous simulations have demonstrated that there is a third type of fluctuating magnetic

277 field near the shock which is not a normal mode of linear theory but corresponds to the

278 magnetosonic-like character of the shock itself, with $|\delta B_{\|}|^2 >> |\delta B_{\perp}|^2$. Such fluctuations

279 are intrinsic to one-dimensional shock simulations [e.g., *Liewer et al.*, 1993; *Kucharek and*

280 *Scholer*, 1995; *Wu et al.*, 2009], and have been demonstrated in the two-dimensional bow

281 shock simulations of *McKean et al.* [1995]. We call these "non-modal" fluctuations, and

282 show that they propagate primarily at $k_{\|} = 0$, although they have smaller amplitude

283 oblique components as well.

284 Magnetic fluctuations from both the Alfvén-cyclotron and proton mirror instabilities

285 are excited downstream of the shock in the present two-dimensional simulation as shown

286 by Figure 3, which displays B_x, $\delta B_y = B_y - \overline{B_y}$ ($\overline{B_y}$ is subtracted away to highlight

287 the remaining smaller amplitude fluctuations in B_y), and B_z at $\Omega_i t = 60$. The Alfvén-

288 cyclotron modes correspond to magnetic fluctuations in B_z because the mirror mode and

289 the non-modal fluctuations have insignificant B_z components. As shown by Figure 3c, the

290 fluctuations in B_z have locally field-aligned wavenumbers. Near the shock, the solar wind

291 ions have a large anisotropy (as shown in Figure 2e) and drive Alfvén-cyclotron modes at

292 relatively short wavelengths; the resulting enhanced fluctuations rapidly reduce the driving

293 anisotropy. Once the free energy is depleted in the thermal or solar wind component,

294 the anisotropy on the pickup ions or suprathermals becomes the dominant driver for

295   the instability and the resulting enhanced fluctuations have relatively longer wavelengths

296   (Compare Figures 11a and 11b). Far downstream of the shock, both anisotropies are

297   relaxed and the remaining Alfvén-cyclotron fluctuations damp with x.

298   In Figure 3a, the fluctuations in B_x, like those in B_z, propagate at $\mathbf{k} \times \mathbf{B}_o = 0$, showing

299   that these fluctuations are predominantly the corresponding coplanar transverse compo-

300   nent of the Alfvén cyclotron fluctuations in B_z. In contrast, the fluctuations in δB_y of

301   Figure 3b are very different from those in B_x and B_z. Figure 3b demonstrates large ampli-

302   tude fluctuations of short scales near the shock. Further downstream, these fluctuations

303   become weaker and demonstrate oblique propagation with longer wavelengths. Compres-

304   sive modes at oblique propagation are characteristic of the proton mirror instability.

305   Figure 4 shows the magnetic fluctuation spectra as functions of k_y at two different x

306   locations and at $\Omega_i t = 60$. The left panels are for magnetic fluctuations at $x = 210\lambda_i$,

307   immediately downstream of the shock. The right panels are for magnetic fluctuations

308   at $x = 300\lambda_i$, far downstream of the shock. From top to bottom, panels (a) and (b)

309   present $|B_x(k_y)|^2$, panels (c) and (d) show $|\delta B_y(k_y)|^2$, and panels (e) and (f) illustrate

310   $|B_z(k_y)|^2$. This figure shows that most wavelength fluctuations decrease with increasing

311   x for all three magnetic components. Furthermore, the maximum fluctuating magnetic

312   energy shifts to longer wavelengths as distance from the shock increases. However, the

313   transverse components, $|\delta B_x(k_y)|^2$ and $|\delta B_z(k_y)|^2$, retain spectral peaks at $k_y > 0$ in the

314   far downstream, indicating that the Alfvén-cyclotron fluctuations have maintained weak

315   growth, whereas the $|\delta B_y(k_y)|^2$ compressive component spectrum decreases montonically

316   with k_y, showing that the mirror-like fluctuations are consistently damped here.

317   Figure 5a displays magnetic fluctuation spectra $|B_z(k_y)|^2$ at $\Omega_i t = 60$ as functions of

318 x. Figure 5b is the same as Figure 5a but has curves superimposed on it to represent

319 the most unstable k_y calculated from linear dispersion theory of the Alfvén-cyclotron

320 instability using the local plasma parameters given by Figure 2 with the assumption that

321 both ion components have bi-Maxwellian velocity distributions. The red and blue curves

322 represent instabilities driven by anisotropies of the solar wind ions and the pickup ions,

323 respectively (See the Appendix). The dashed part of each curve means that the growth

324 rate of the corresponding mode is smaller than that of the other mode. The most intense

325 wavenumbers from the simulation agree with the linear theory prediction reasonably well,

326 confirming our interpretation that the enhanced fluctuations in B_z are generated by the

327 Alfvén-cyclotron instability.

328   Figure 6 presents a two-dimensional spectral analysis of the magnetic fluctuations over

329 $200\lambda_i < x < 300\lambda_i$ at $\Omega_i t = 60$. Panels a), b), c), and d) display spectra of B_x, B_y (with

330 the areal average of B_y subtracted away), B_z and $\delta B_y = B_y - \overline{B_y}$, respectively. Here

331 $|B_x(k_x, k_y)|^2$ and $|B_z(k_x, k_y)|^2$ are similar and further confirm the field-aligned propaga-

332 tion feature of these fluctuations. The dominant fluctuations clearly have $k_x = 0$ and

333 $0.15 < k_y\lambda_i < 0.50$, consistent with the linear dispersion theory prediction for Alfvén-

334 cyclotron waves as shown in Figure 5b. The spectrum of B_y in Figure 6b has dominating

335 fluctuations with $k_y = 0$, which are the features of non-modal fluctuations demonstrated

336 by previous one-dimensional simulations of perpendicular and quasi- perpendicular shocks.

337 But there are also relatively weak oblique fluctuations with nonzero k_y $(k_y\lambda_i \approx 0.1)$ in

338 $|B_y(k_x, k_y)|^2$, which do not show up very well due to the dominance of the $k_y = 0$ fluc-

339 tuations. These oblique fluctuations stand out in Figure 6d after the $k_y = 0$ fluctuations

340 are removed subtracting $\overline{B_y}$ from B_y before the spectral analysis. As discussed below,

341 these oblique fluctuations are a combination of the mirror modes and some oblique non-

342 modal fluctuations. From $\nabla \cdot \mathbf{B} = 0$, there should be B_x fluctuations in the present

343 two-dimensional simulation associated with these oblique fluctuations in B_y. This is in-

344 deed the case as shown by the oblique fluctuation in Figure 6a, which is at the same

345 location ($k_x \lambda_i \approx 0.1$, $k_y \lambda_i \approx 0.08$) as the strongest fluctuation in Figure 6d.

346   In order to identify the oblique fluctuations in Figure 6d, the correlation coefficient be-

347 tween $\delta B_y = B_y - \overline{B_y}$ and the density fluctuations $\delta n_i = n_i - \overline{n_i}$ [similar to $\overline{B_y}$, $\overline{n_i}$ is $n_i(x, y)$

348 averaged over y] as a function of x is displayed in Figure 7. The density fluctuations are

349 rather noisy due to limited number of particles per cell, so this coefficient exhibits rather

350 large amplitude oscillations. Nevertheless, it is clear that the correlation coefficient is

351 close to 1 near the shock, decreases rapidly thereafter, and displays small but negative

352 values farther downstream of the shock. Near the shock, the density fluctuations and

353 the magnetic fluctuations are positively correlated, indicative of the basic magnetosonic

354 character of the shock. Farther downstream, the negative correlation between the den-

355 sity fluctuations and the magnetic fluctuations further confirms out interpretation that

356 mirror-like fluctuations are present here.

357   Figure 8 displays magnetic fluctuation component energy densities $\overline{B_x^2}$ (red curve), $\overline{\delta B_y^2}$

358 (black curve), and $\overline{B_z^2}$ (blue curve) versus x. The energy in B_z, which indicates the devel-

359 opment of Alfvén-cyclotron modes, grows quickly after the shock, reaches saturation and

360 then decreases farther downstream. This is consistent with the picture that these fluctu-

361 ations are generated through the Alfvén-cyclotron instability driven by the anisotropies

362 in both ion components as shown in Figure 2 and as indicated by the linear theory of

the Appendix. The first maximum at $x = 200\lambda_i$ in $\overline{B_z^2}$ is due to the short-wavelength

modes driven by the solar wind ion anisotropy. These modes decrease thereafter while the

longer wavelength modes driven by the pickup ions still increase and reach a maximum

amplitude at $x = 215\lambda_i$. The $\overline{B_x^2}$ demonstrates similar behavior as $\overline{B_z^2}$ suggesting once

again that the contributions of the mirror modes and the non-modal fluctuations in B_x

are small compared with those of the Alfvén-cyclotron mode.

However, the compressive component of the magnetic spectrum behaves differently than

the other two and reaches a large maximum right at the shock overshoot. Our interpre-

tation is that this maximum is due to the non-modal fluctuations as suggested by Figure

7. The oblique non-modal fluctuations decay quickly as one goes downstream, but have

strong coupling to the mirror-like fluctuations at the shock, which persist deeper into the

downstream. This is different from the Alfvén-cyclotron fluctuations, which have to grow

from relatively small amplitude initial perturbations in B_x and B_z, and explains why the

amplitude of the mirror mode fluctuations are relatively large compared to what might

be expected from predictions of linear dispersion theory. The energy of the mirror mode

fluctuations reaches its maximum at $215\lambda_i$ as well, where the large pickup ion anisotropy

ceases. After that, both the Alfvén-cyclotron and mirror mode fluctuations decrease

while the Alfvén-cyclotron fluctuations damp slightly slower, consistent with the linear

dispersion theory prediction that the Alfvén cyclotron instability has a smaller anisotropy

threshold than the mirror instability given the local plasma parameters from Figure 2.

At $x = 300\lambda_i$, the energy of the Alfvén-cyclotron fluctuations is $\overline{B_x^2} + \overline{B_z^2} \approx 0.2$ while

the energy of the mirror mode fluctuations is $\leq \overline{\delta B_y^2} \approx 0.04$. Thus the Alfvén-cyclotron

mode dominates the mirror mode sufficiently far downstream from the shock, although

386 the mirror mode is much stronger than one might expect from the predictions of linear

387 dispersion theory.

388 The primary consequence of wave-particle interactions by the dominant Alfvén-

389 cyclotron fluctuations is ion pitch-angle scattering, so that the high-speed part of

390 suprathermal ion perpendicular velocity distributions in the heliosheath should be de-

391 pleted by comparison with those computed from one-dimensional hybrid simulations.

392 This is demonstrated in Figure 9, which compares the downstream perpendicular ion

393 velocity distributions at $\Omega_i t = 60$ from one-dimensional and two-dimensional simulations

394 with the same upstream conditions. Both distributions are obtained by directly binning

395 ions in their respective speed space. This figure shows clearly that suprathermal ions at

396 $v_\perp > 10 v_A$ are depleted, whereas suprathermals at $v_\perp \leq 10 v_A$ are enhanced, consistent

397 with the physical picture of ion pitch-angle scattering by Alfvén-cyclotron fluctuations.

6. Summary and Conclusions

398 This paper has described a two-dimensional, zero-electron-mass hybrid simulation of

399 the perpendicular termination shock. This simulation generalizes the one-dimensional

400 computations of *Wu et al.* [2009, 2010], and demonstrates properties of the enhanced

401 fluctuations which arise in the heliosheath.

402 Analysis of our simulation shows three distinct types of fluctuations at and downstream

403 of the shock: Alfvén-cyclotron modes, mirror-like modes, and "non-modal" fluctuations.

404 The first two of these correspond to normal modes which can be described by linear disper-

405 sion theory in the relatively uniform downstream plasma, whereas the non-modal fluctua-

406 tions are intrinsically nonlinear oscillations associated with the strong inhomogeneities at

407 and near the shock. The Alfvén-cyclotron fluctuations have magnetic fluctuations primar-

ily perpendicular to the background magnetic field $\mathbf{B}_o$, $|\delta B_\perp|^2$, whereas both the mirror

fluctuations and the non-modal fluctuations are strongly compressive, with $|\delta B_\parallel|^2$ dom-

inating the magnetic fluctuation energy. Mirror-like fluctuations have anticorrelated δn

and $\delta B_\parallel$, whereas the magnetosonic-like non-modal fluctuations have positively correlated

density and compressive magnetic fluctuations.

We interpret the downstream enhanced fluctuations in our simulation as follows. The

primary driver of the mirror-like fluctuations is their coupling to the compressive mag-

netic component of the non-modal fluctuations near the shock. Thus Figure 8 shows an

almost-monotonic decay of $|\delta B_\parallel|^2$ as x increases downstream away from the shock. The

Alfvén-cyclotron downstream fluctuations are driven partially by coupling to the weaker

transverse magnetic component of the non-modal fluctuations, but more strongly through

instabilities excited by the ion component temperature anisotropies at the shock (as il-

lustrated in Figure 2). Immediately downstream of the shock, the instability is driven

by the thermal component $T_{\perp t}/T_{\parallel t} >> 1$ (Figure 10a); the corresponding wavelength is

quite short (Figure 11a). After rapid reduction of the thermal component anisotropy,

the weaker $T_{\perp s}/T_{\parallel s} > 1$ remains and drives a weaker instability (Figure 10b) at a longer

wavelength (Figure 11b). Thus the relatively constant magnitude of $|\delta B_{\perp\perp}|^2$ over 200

$\lesssim x/\lambda_i \lesssim$ 250 in Figure 8 indicates a competition between the decay of modes driven

at the shock and the growth of fluctuations driven by the Alfvén-cyclotron instability.

Well downstream of the shock, our simulation further shows that Alfvén-cyclotron fluc-

tuations have substantially larger amplitudes than mirror-like fluctuations. This is con-

sistent with low-Mach low-β bow shock results, both simulated [*McKean et al.*, 1996]

and observed [*Sckopke et al.*, 1990]. However, high Alfvén-Mach bow shock results, both

431 simulated *McKean et al.* [1995] and observed [*Sckopke et al.*, 1990], show that transverse

432 and compressive magnetic fluctuations have similar downstream amplitudes.

433 We here offer a hypothesis to explain these results. The intrinsic non-modal fluctuations

434 at perpendicular and quasi-perpendicular shocks have both a strong $\delta B_\parallel$ component and

435 a significant $k_\perp$ component. Therefore, they couple strongly to mirror-like fluctuations

436 (which also have significant $\delta B_\parallel$ and $k_\perp$ components) and directly excite those modes

437 via nonlinear interactions. The Alfvén-cyclotron fluctuations with $\delta B_\parallel = 0$ and $k_\perp =$

438 0 do not couple as strongly to the non-modal oscillations; they are more likely to arise

439 via the development of $T_{\perp i}/T_{\parallel i} > 1$ at the shock and the subsequent linear growth of the

440 Alfvén-cyclotron instabilities downstream. Our hypothesis is that the nonlinear excitation

441 of mirror-like fluctuations increases more rapidly with magnetosonic Mach number than

442 does the linear growth of the Alfvén-cyclotron instability. Thus the high Alfvén-Mach

443 bow shock simulations of *McKean et al.* [1995] shows the two modes develop similar

444 amplitudes downstream, whereas Alfvén-cyclotron modes dominate low Alfvén-Mach bow

445 shock simulations [*McKean et al.*, 1996]. Although our termination shock simulation is

446 run at $M_A = 8$, it is of course the magnetosonic Mach number which determines the

447 strength of a perpendicular shock; as $M_{ms} < 4$ here due to the high-β contribution of

448 the pickup ions, our hypothesis is consistent with Alfvén-cyclotron modes dominating the

449 downstream fluctuations. Further two-dimensional simulations of the termination shock

450 at various Alfvén-Mach numbers and with various pickup ion upstream densities will be

451 needed to test this idea more fully.

452 Thus we conclude that the enhanced fluctuations of the near-shock heliosheath should

453 consist primarily of Alfvén-cyclotron modes. We further conclude that the consequences

454 of wave-particle interactions in this region should be primarily pitch-angle scattering of

455 relatively fast suprathermal ions, as indicated by Figure 9.

456   Observations of the heliosheath magnetic field by Voyager 1 and 2 [*Burlaga et al, 2005,*

457 *2008*] show that the field fluctuations tend to be large amplitude, widely varying and

458 very complex. There are indications of the presence of magnetic holes [*Burlaga et al.,*

459 *2007*] which could be signatures of mirror waves [*Winterhalter et al.*, 1994; *Zhang et al.,*

460 2008], but identification of individual wave modes in the Voyager data is very difficult.

461 Instead, statistical methods have been used to characterize the distribution of magnetic

462 field amplitudes in the heliosheath [*Burlaga et al., 2009*]. Detailed measurements of the

463 magnetic field strength as Voyager 2 crossed the termination shock (e. g., the crossing

464 termed TS-3) [*Burlaga et al., 2008*] and entered the heliosheath show the jump in the

465 field at the shock and the fluctuations just downstream of the shock, which are generally

466 consistent with the simulation results in Figure 1, but corresponding wave spectra are not

467 available for more detailed comparison. The Alfvén-cyclotron and mirror modes expected

468 downstream of the shock scatter ions differently, and their relative amplitudes in the

469 heliosheath affect the consequent ENA fluxes which emanate from this region. However,

470 a detailed comparison of the scattering consequences of these two modes is beyond the

471 purview of this manuscript.

472   We emphasize that our conclusions are based upon a two-dimensional simulation model

473 of the perpendicular termination shock. Two natural extensions of this work would be

474 to consider the more general quasi-perpendicular shocks and fully three-dimensional sim-

475 ulations. In particular the three-dimensional computations of *Shoji et al.* [2009] suggest

476 that compressible mirror-like fluctuations may persist deeper into the downstream than

477 the incompressible Alfvén-cyclotron fluctuations.

Appendix A: Linear Theory

478 This Appendix describes results obtained from the kinetic linear dispersion equation

479 for waves and instabilities in a homogeneous, collisionless, magnetized plasma [*Gary*,

480 1993]. We here use linear theory to compute growth rates for instabilities driven by the

481 anisotropy $T_{\perp j}/T_{\parallel j} > 1$ on the thermal and suprathermal ion components downstream of

482 the termination shock. The theory is based upon the assumption that both components

483 can be represented as bi-Maxwellian velocity distributions. Observations [*Funsten et al.*,

484 2009], models [*Zank et al.*, 2010 and references therein], and termination shock simulations

485 [*Wu et al.*, 2010] all indicate that suprathermal ion velocity distributions may have power-

486 law velocity dependence at relatively large speeds. Nevertheless, we argue that, because

487 we are computing the growth rates of strongly resonant kinetic instabilities, it is the

488 properties of the mid-velocities of a component distribution, not the high-velocity tails,

489 which are most important for determining instability properties, in which case the use of

490 bi-Maxwellian distributions is likely to be an appropriate approximation.

491 We use parameters from the hybrid simulations described in section 4 [See also Table

492 2 of *Wu et al.*, 2009] for two cases corresponding to immediately downstream and well

493 downstream of the shock. In both cases the proton component anisotropies lead to the

494 growth of two distinct instabilities: the Alfvén-cyclotron instability with maximum growth

495 at $\mathbf{k} \times \mathbf{B}_o = 0$ and $0 < \omega_r < \Omega_p$, and the proton mirror instability with $\omega_r = 0$ and

496 maximum growth rate at propagation oblique to $\mathbf{B}_o$.

For a strictly perpendicular shock, there is no parallel heating of the plasma, so immediately downstream we take $\beta_{\|t} = 0.04$ and estimate $T_{\|s}/T_{\|t} = 800.0$ as well as $T_{\perp s}/T_{\|s} = 2.5$, consistent with the results of the simulations described in Section 4. We assume the electrons are isotropic and $T_{\|e}/T_{\|t} = 1.0$. With these parameters, we compute the maximum growth rate γ_m/Ω_p and the corresponding wavenumber $k_m c/\omega_i$ and angle of propagation θ_m as functions of the thermal component anisotropy, with results illustrated in Figures 10a through 12a.

If the thermal component anisotropy is relatively weak, both instabilities are driven by the suprathermal anisotropy; for these parameters the Alfvén-cyclotron instability has the consistently larger value of γ_m/Ω_p. As $T_{\perp t}/T_{\|t}$ increases, the growth rates of the modes driven by this thermal component anisotropy also increase, and their contributions appear on the right-hand sides of the figures. For our parameters, the Alfvén-cyclotron suprathermal instability has maximum growth rate at $0.25 < k_\| c/\omega_p < 0.30$, but the Alfvén-cyclotron thermal mode has maximum growth at $k_\| c/\omega_p > 1$. Similarly, the suprathermal mirror instability has γ_m at $kc/\omega_p \simeq 0.185$ while the thermal mirror mode has maximum growth at $kc/\omega_p > 1$. Figure 12a shows that θ_m of the suprathermal mirror instability is a weak function of the thermal proton anisotropy, but that the thermal mirror instability propagates at successively steeper angles with respect to $\mathbf{B}_o$ as $T_{\perp t}/T_{\|t}$ decreases and the instability is correspondingly weakened.

The simulations demonstrate that a short distance downstream of the shock, the thermal protons are scattered to near-isotropy. For our calculations in this second regime, we assume $T_{\perp t}/T_{\|t} = 1.5$, $T_{\|s}/T_{\|t} = 30.0$, and $\beta_{\|t} = 0.80$. Using these parameters, we plot in Figure 10b the maximum growth rates of the Alfvén-cyclotron and mirror insta-

520 bilities driven by the suprathermal component anisotropy as a function of $T_{\perp s}/T_{\parallel s}$. Once

521 again, the Alfvén-cyclotron mode has the consistently higher growth rate, so that we

522 expect this mode to dominate magnetic field fluctuations in the far downstream, in agree-

523 ment with the simulations described above. Figure 11b shows that the Alfvén-cyclotron

524 instability here has a relatively long wavelength, with $0.10 < k_{\parallel}c/\omega_p < 0.30$ at maxi-

525 mum growth. Figure 12b shows that the angle of propagation at maximum growth rate

526 of the suprathermal mirror instability becomes more nearly perpendicular to $\mathbf{B}_o$ as the

527 suprathermal anisotropy decreases and the instability weakens.

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Figure 1. B_y profiles at simulation time $\Omega_i t = 60$ and different y-values (black lines): (a) $y = 50\lambda_i$, (b) $y = 125\lambda_i$, and (c) $y = 200\lambda_i$. The gray line in each panel is B_y averaged over y ($\overline{B_y}$).

Figure 2. Plasma properties of the solar wind ions (left column) and pickup ions (right column) as functions of x at simulation time $\Omega_i t = 60$. The panels present, in succession from top to bottom, the ion component density averaged over y and normalized to the upstream total plasma density, the ion component temperatures normalized to the average upstream value of $m_i v_A^2$ where $T_j = m_i < v_j^2 - < v_j >^2 >$ and $<>$ represents a velocity average over the ion component j (T_x, T_y, and T_z are represented by red, black, and blue lines, respectively), the component temperature ratio $T_{\perp j}/T_{\parallel j} = (T_{xj} + T_{zj})/(T_{yj}/2)$, and the ion component $\beta_{\parallel j}$.

Figure 3. The magnetic fluctuations excited downstream of the shock at simulation time $\Omega_i t = 60$: (a) $B_x(x,y)$, (b) $\delta B_y(x,y) = B_y(x,y) - \overline{B_y(x,y)}$ and (c) $B_z(x,y)$.

Figure 4.   Spectra of magnetic fluctuations as functions of k_y at two x locations and at simulation time $\Omega_i t = 60$. The left panels are at $x = 210\lambda_i$, immediately downstream of the shock. The right panels are at $x = 300\lambda_i$, far downstream of the shock. The panels from top to bottom present the spectra of $|B_x(k_y)|^2$ (panels a and b), $|\delta B_y(k_y)|^2$ (panels c and d), and $|B_z(k_y)|^2$ (panels e and f), respectively.

Figure 5.   (a) The spectra $|B_z(k_y)|^2$ at simulation time $\Omega_i t = 60$ as a function of x. (b) Same as (a), but with curves superimposed to represent the most unstable k_y calculated from linear dispersion theory of the Alfvén-cyclotron instability using the local plasma parameters given by Figure 2. The red curve and the blue curve are for modes driven by the anisotropies of the solar wind and pickup ions, respectively. The dashed part of each curve means that the growth rate of the corresponding mode is smaller than the growth rate of the other mode.

Figure 6.   Two-dimensional wavenumber spectra of the three components of magnetic fluctuations over $200\lambda_i < x < 300\lambda_i$ at simulation time $\Omega_i t = 60$: (a) $|B_x(k_x, k_y)|^2$, (b) $|B_y(k_x, k_y)|^2$ (with the areal average of B_y subtracted away), (c) $|B_z(k_x, k_y)|^2$, and (d) $|\delta B_y(k_x, k_y)|^2$.

Figure 7.   The correlation coefficient between $\delta B_y = B_y - \overline{B_y}$ and density fluctuations $\delta n_i = n_i - \overline{n_i}$ as a function of x.

Figure 8. Magnetic fluctuation component energies as functions of x at simulation time $\Omega_i t = 60$. Here the red curve indicates $\overline{B_x^2}$, the black curve indicates $\overline{\delta B_y^2}$ and the blue curve denotes $\overline{B_x^2}$.

Figure 9. Ion perpendicular velocity distributions $v_\perp f(v_\perp)$ downstream of the shock summed over all y values and $250 \leq x/\lambda_i \leq 350$ at simulation time $\Omega_i t = 60$ from a one-dimensional simulation (dashed line) and from the two-dimensional simulation (solid line).

Figure 10. Linear theory results: The maximum growth rates of Alfvén-cyclotron (solid lines) and proton mirror (dashed lines) instabilities from kinetic linear dispersion theory. For both panels, $n_s/n_o = 0.20$. (a) Maximum growth rates as a function of the thermal proton component anisotropy. Parameters here are characteristic of those immediately behind a termination shock: $\beta_{\|t} = 0.04$, $T_{\|s}/T_{\|t} = 800.0$ and $T_{\perp s}/T_{\|s} = 2.5$. The left-hand, relatively horizontal lines represent instabilities driven by the suprathermal ion anisotropy, whereas the right-hand, upward trending lines represent instabilities driven by the thermal ion anistropy. (b) Maximum growth rates of instabilities driven by the suprathermal proton anisotropy, as functions of that anisotropy. Parameters here are characteristic of those somewhat downstream of a termination shock: $\beta_{\|t} = 0.80$, $T_{\|s}/T_{\|t} = 30.0$ and $T_{\perp t}/T_{\|t} = 1.5$.

Figure 11.   Linear theory results: The dimensionless wavenumbers at maximum growth rates of the Alfvén-cyclotron (solid lines) and proton mirror (dashed lines) instabilities from kinetic linear dispersion theory. Parameters are the same as for the corresponding panels in Figure 10. (a) Wavenumbers as functions of the thermal proton anisotropy. The upper two curves correspond to instabilities driven by the thermal proton anisotropy, whereas the lower two curves represent wavenumbers of instabilities driven by the suprathermal proton anisotropy. (b) Wavenumbers of the instabilities driven by the suprathermal proton anisotropy as functions of that anisotropy.

Figure 12.   Linear theory results: The angle at maximum growth rate θ_m of the proton mirror instabilities from kinetic linear dispersion theory. Parameters are the same as for the corresponding panels in Figure 10. (a) Values of θ_m as functions of the thermal proton anisotropy, corresponding to the two different types of mirror instability as labeled. (b) Values of θ_m of the suprathermal mirror instability as a function of the suprathermal anisotropy.

Figure 1

Solar Wind Ions

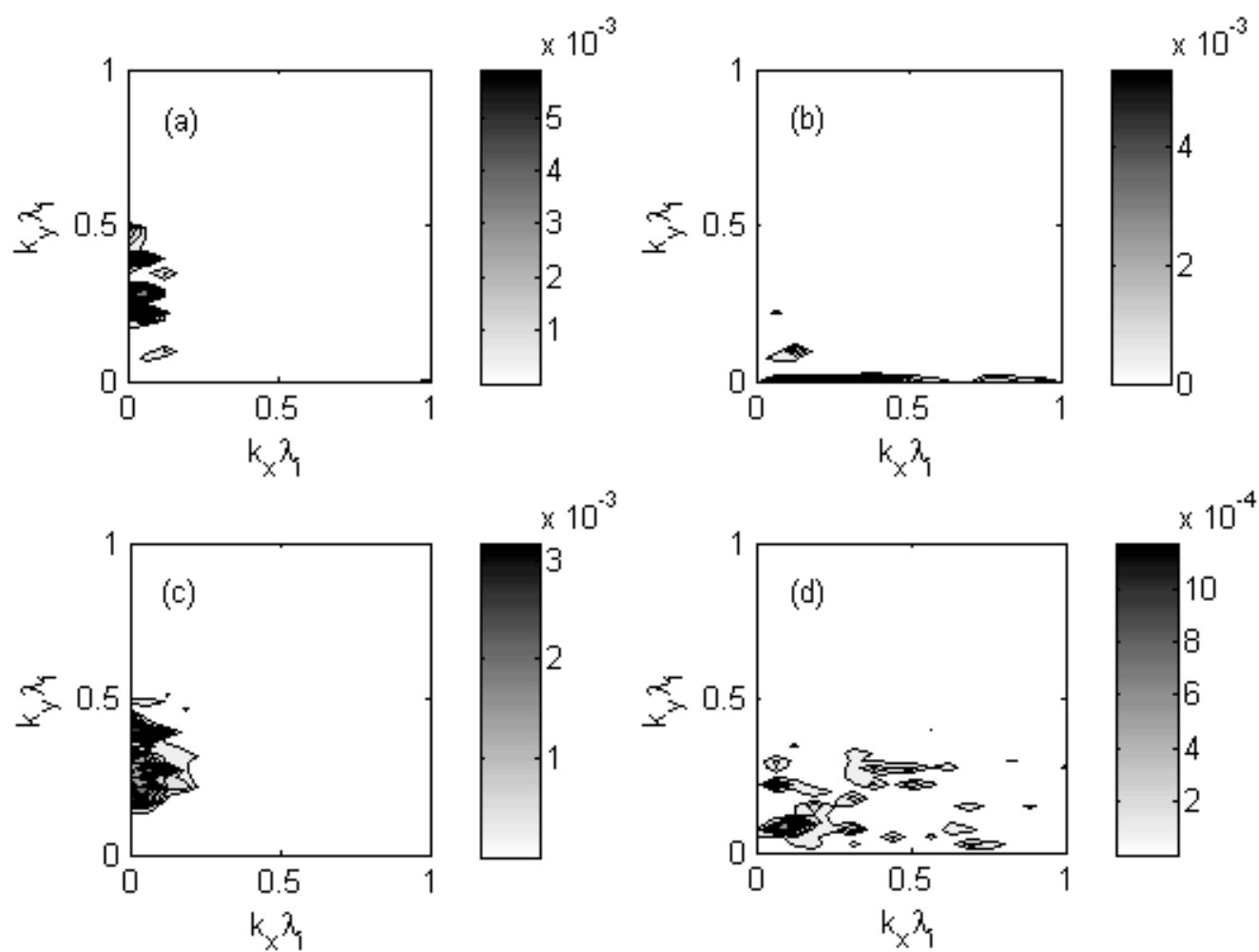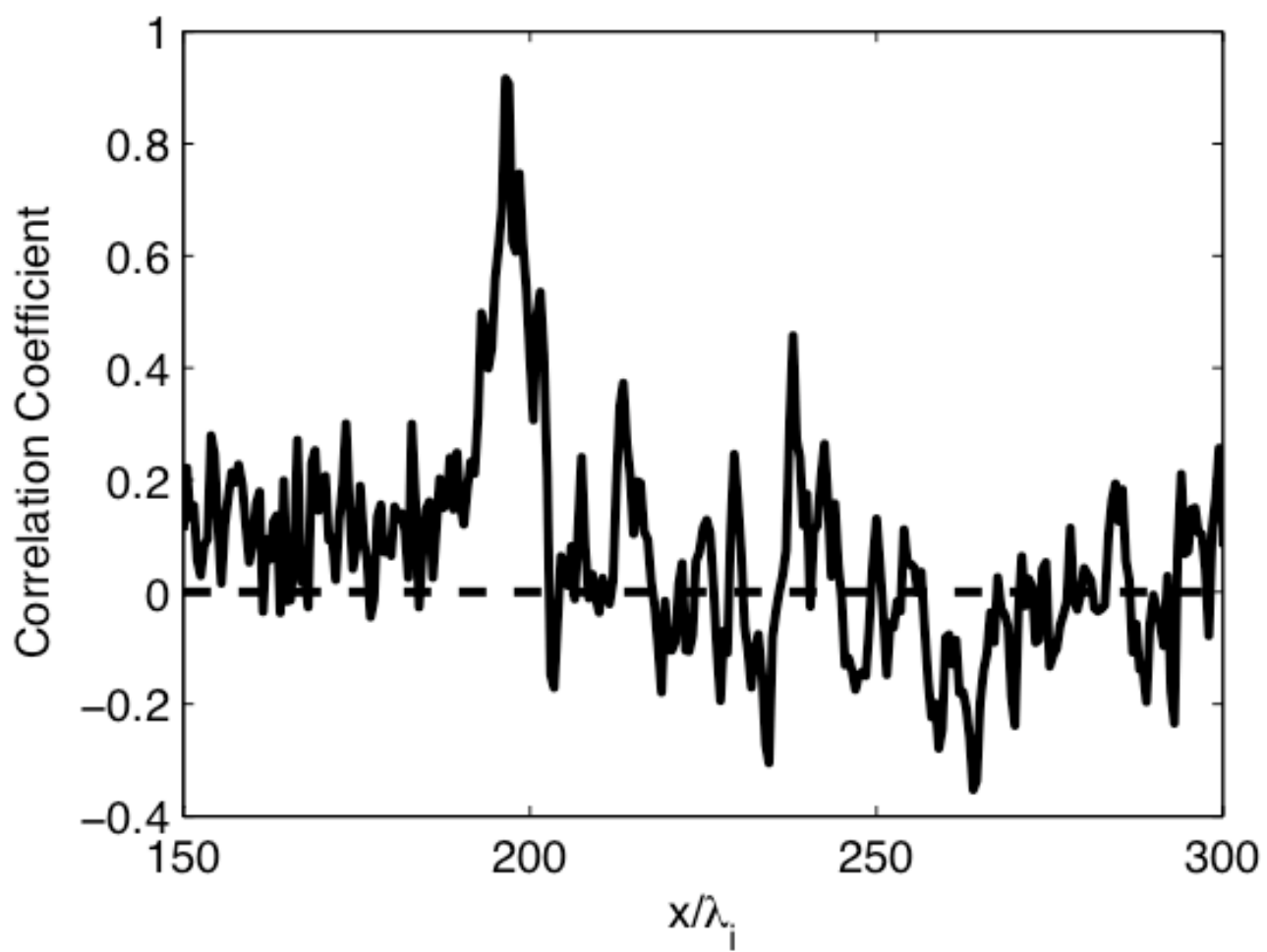Pickup Ions

(a)

(b)

(c)

(d)

(e)
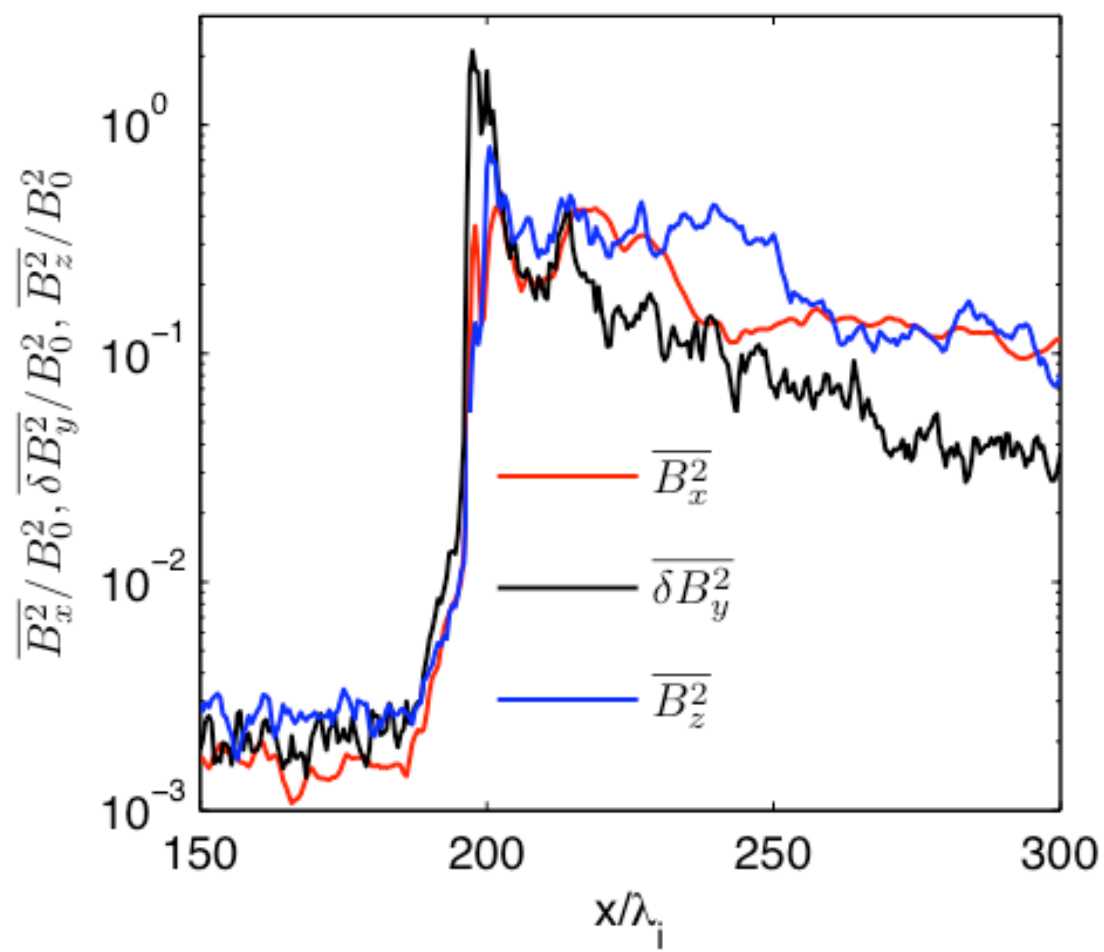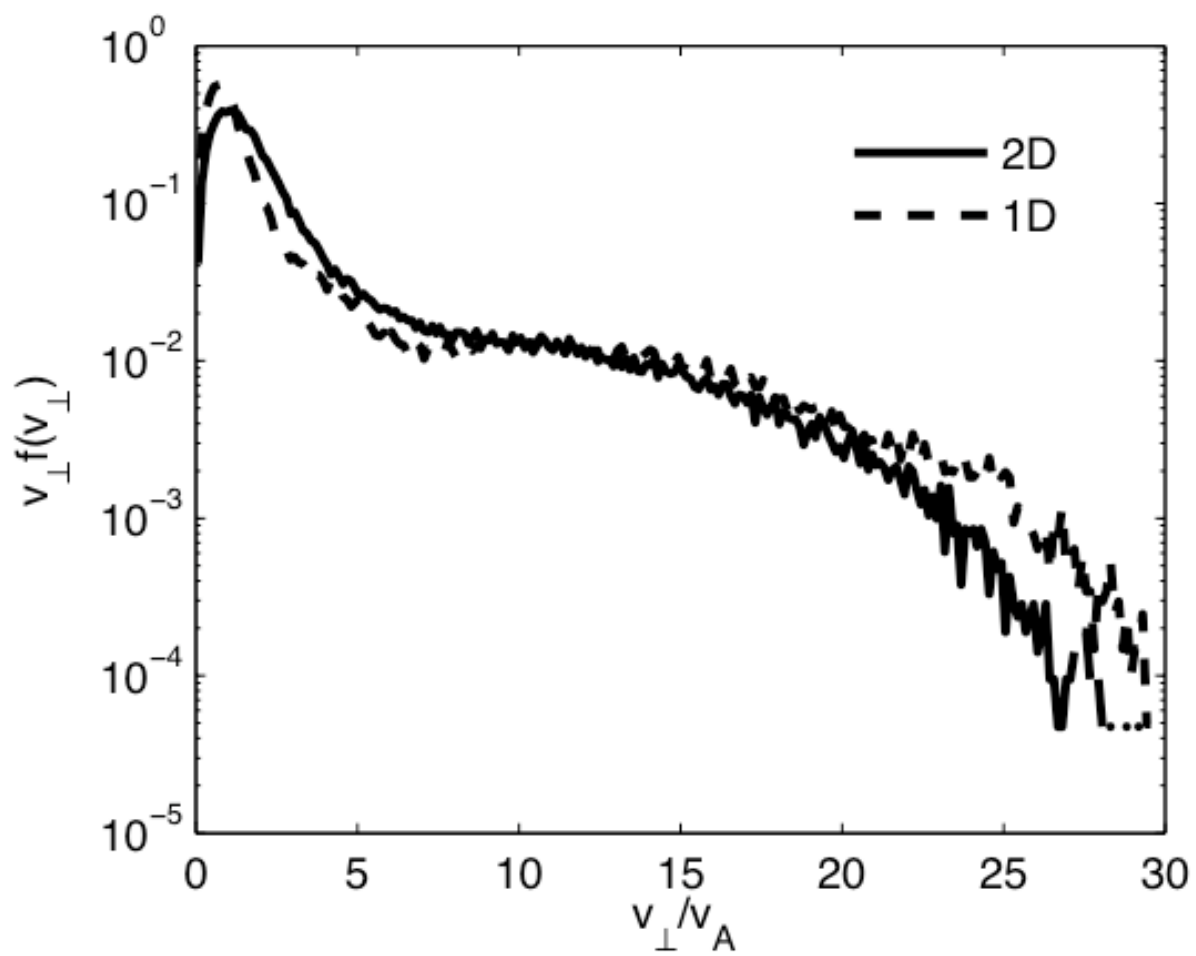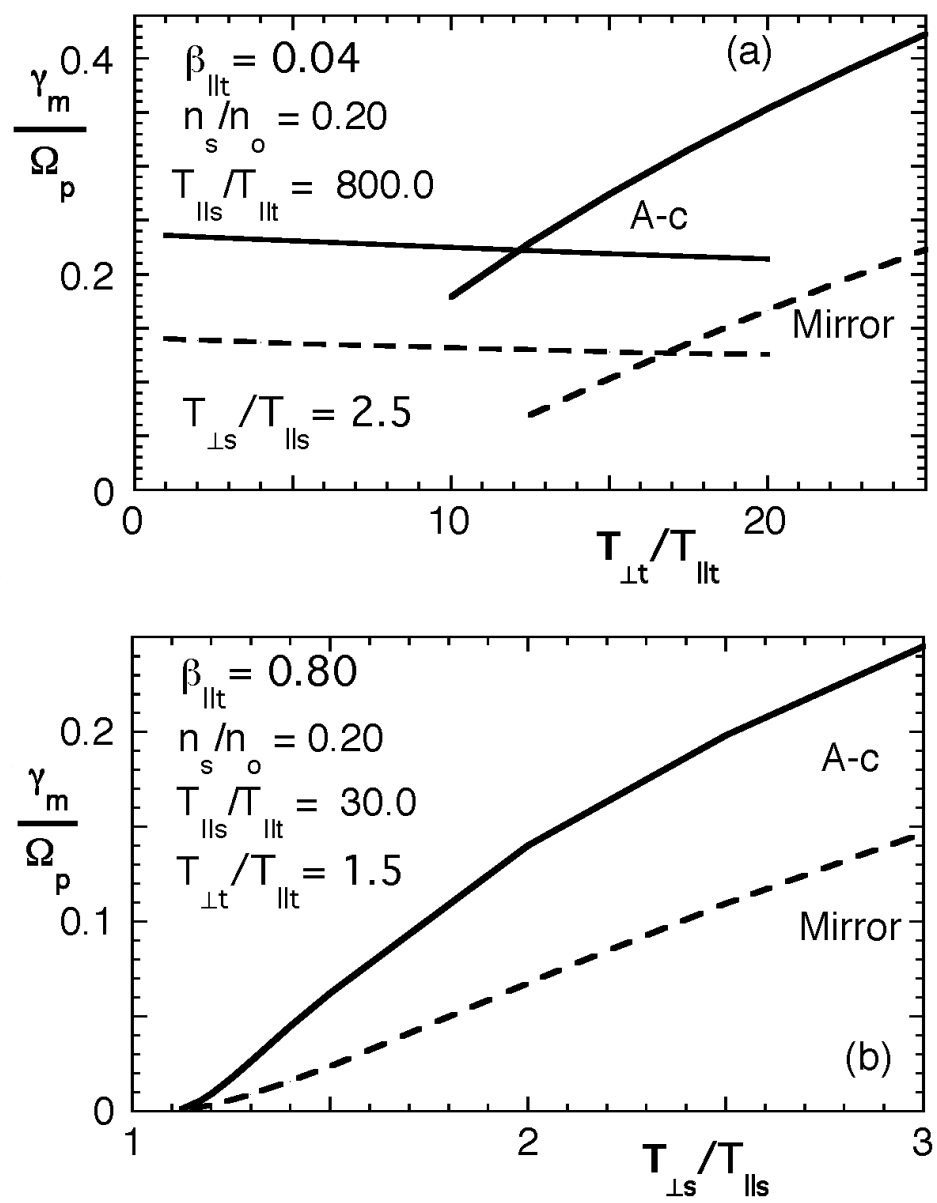
(f)

(g)

(h)
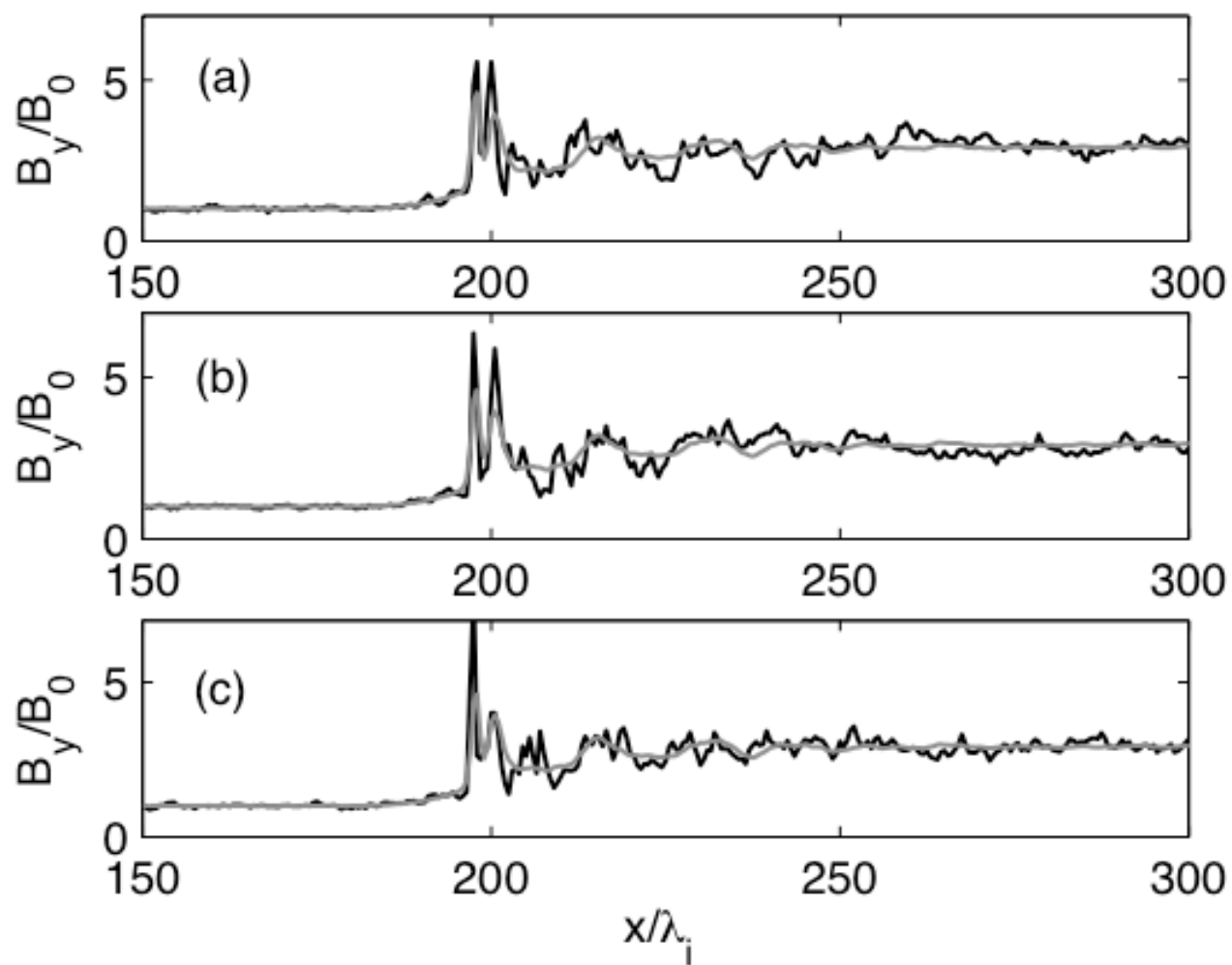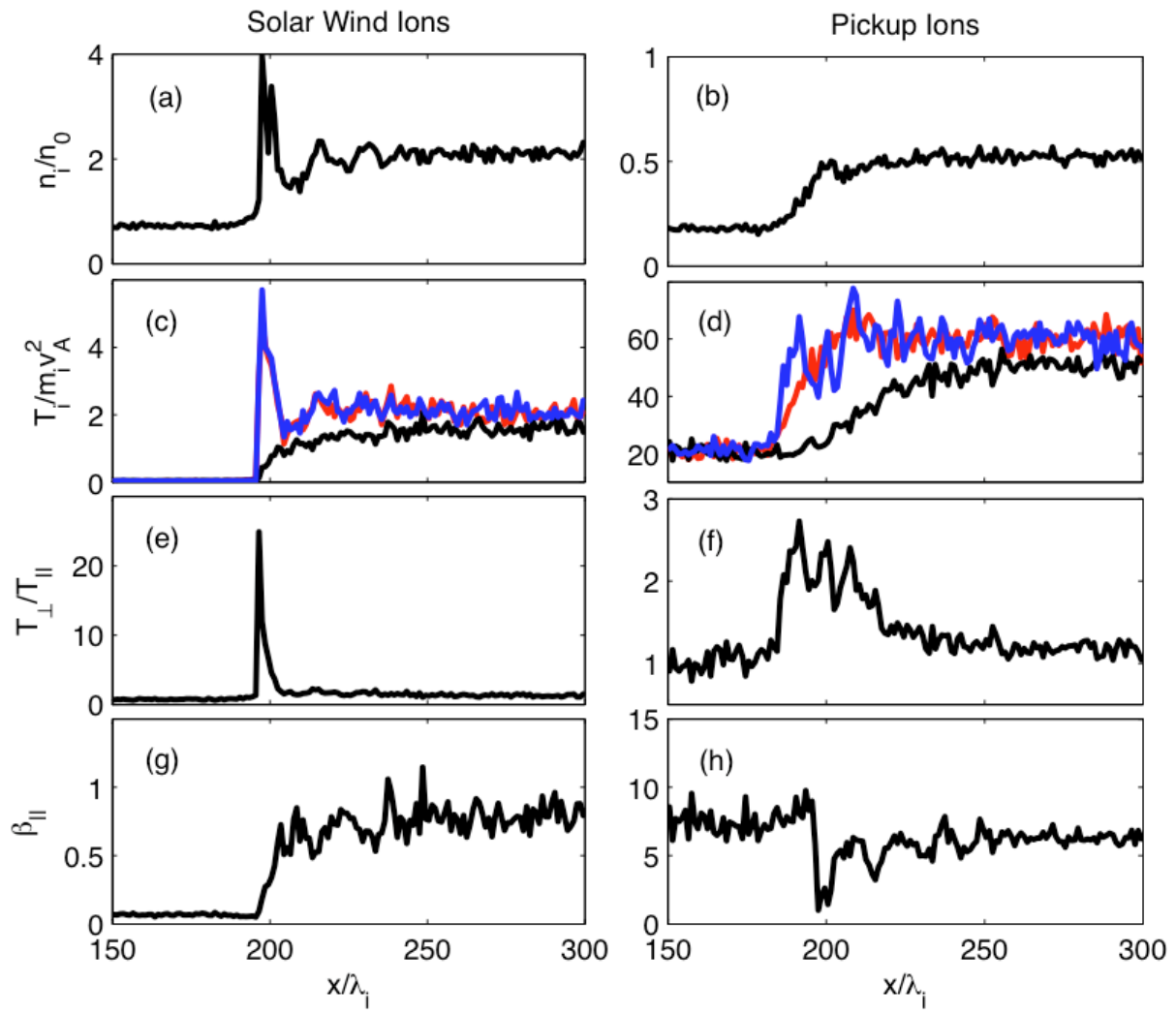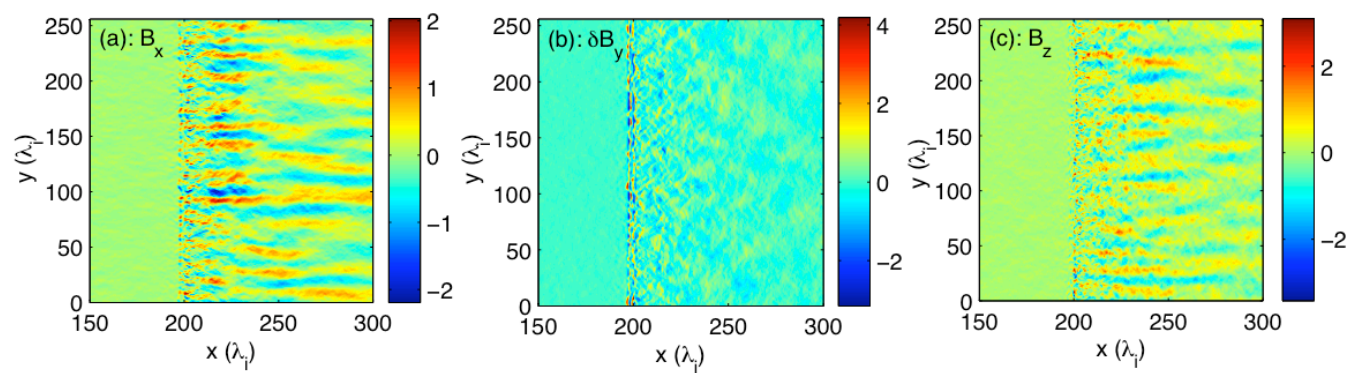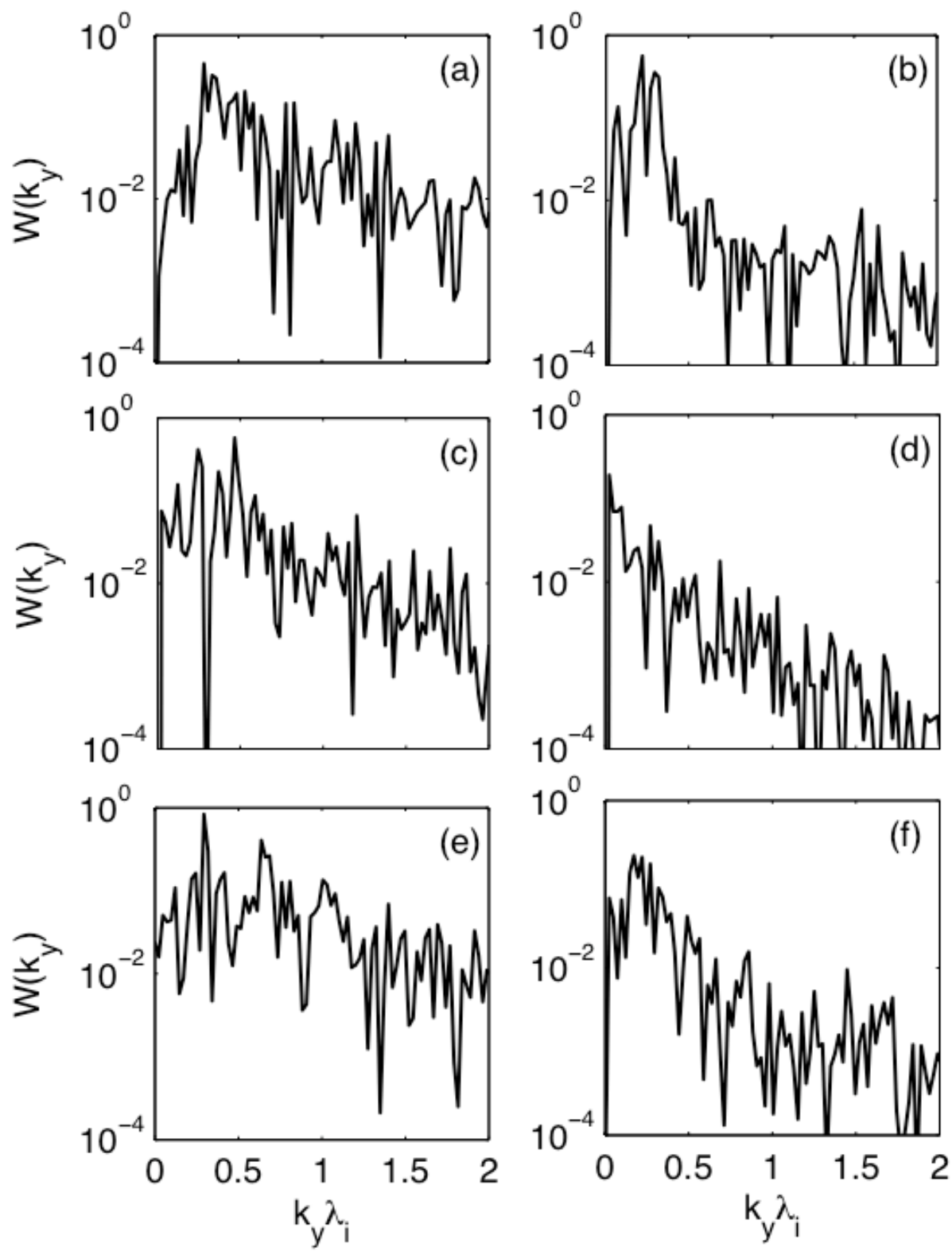
Figure 2

Figure 3

Figure 4
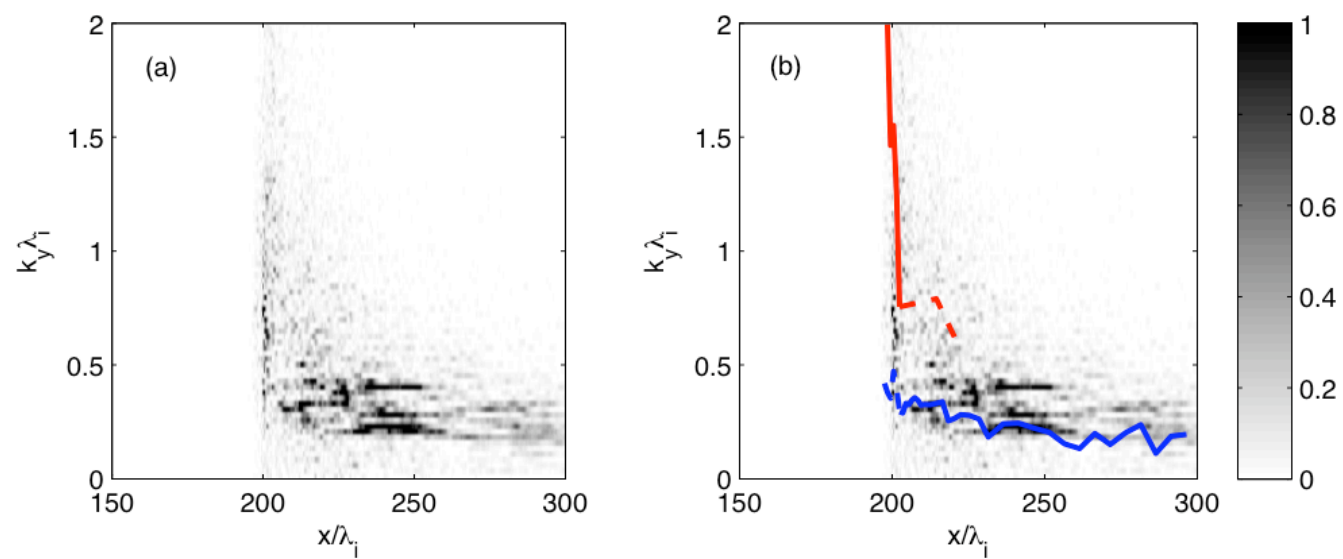
Figure 5

Figure 6

Figure 7

Figure 8

Figure 9

Figure 10

Figure 11

Figure 12